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Did one observe couplings of right-handed quarks to W ?

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I consider non standard EW couplings of light right-handed quarks to W and Z suggested in a systematic non decoupling bottom-up low-energy effective theory approach to possible extensions of the Standard Model. New experimental tests in $K^L_{\mu 3}$ decays based on recent measurements and scalar form factor analysis are discussed. A successful NLO fit to the standard set of Z-pole and other NC data is presented as well.

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1. Introduction

In the Standard Model (SM), right handed fermions do not couple to W and their couplings to Z are proportional to the electric charge. Compelling tests of this feature exist for leptons, whereas for quarks available tests are less conclusive due to the interference with non perturbative QCD effects. Another characteristics of the right-handed sector of the SM is a rather complicated and apriori unexplained spectrum of weak hypercharges. (Since the seventies, the latter has motivated left-right symmetric extensions of the SM [1] which shed a new light on the EW couplings of right handed fermions.) None of the above features of the SM follow from the EW symmetry $S_{EW} = SU(2)_W \times U(1)_Y$, as long as the latter is spontaneously broken: Indeed, with the help of agents of Symmetry Breaking (Higgs fields), it is possible to construct S_{EW} invariant couplings of right handed fermions to W. This fact suggests to look for eventual modifications of the right-handed couplings as a conceivable signal of a non standard EW symmetry breaking. Model independent tests of EWSB require first of all a "bottom-up" Effective Theory approach which starts from the known vertices of the SM and step by step in a low-energy expansion controlled by a **power counting** orders possible non standard effects according to their importance at low energies. Next, it should be specified how the lepton - quark universality could be naturally broken at subleading orders to escape strong experimental constraints concerning leptons.

Such a class of LEETs has been proposed three years ago [2] and further developed and completed later [3]. In this talk, I am going to review the characteristic feature of this class: The appearance at NLO of couplings of right handed quarks to W and modification of their couplings to Z. Then I will comment on first attempts to confront these predictions with experiment [4].

2. Not quite Decoupling EW Low Energy Effective Theory (LEET)

In its minimal version, the LEET contains the naturally light particles of the SM: SU(2) x U(1) gauge fields, chiral fermions (including right-handed neutrinos) and the triplet of GBs. For small momenta $p \ll 4\pi F_W = \Lambda_W \sim 3TeV$, the effective Lagrangian is written as a low-energy expansion

$$\mathscr{L}_{eff} = \sum_{d\geq 2} \mathscr{L}_d, \quad \mathscr{L}_d = \mathscr{O}([p/\Lambda_W]^d) \ , \tag{2.1}$$

where the infrared dimension of a local operator, $d = n_\delta + n_g + n_f/2$ is given by the number of derivatives, the number of gauge couplings and the number of fermion fields. A Feynman diagram with effective vertices v=1... and with L loops counts at low-energy as O(p^d), where

$$d = 2 + 2L + \sum_v (d_v - 2) \ . \tag{2.2}$$

The LEET is renormalizable order by order in the LE expansion, provided at each order, all terms allowed by symmetries are effectively included in (2.1). In particular, the symmetry of the LEET $S_{nat} \supset S_{EW}$ must prevent all "unwanted " non standard vertices to appear already at the leading order $\mathscr{O}(p^2)$. In a bottom-up approach, the higher symmetry S_{nat} is unknown apriori (it is the remnant of the not quite decoupled high energy sector of the theory), but it can be inferred requiring that

the leading order $\mathscr{L}_2$ of the LEET coincides with the Higgs-free part of the SM Lagrangian. I refer to [3], where it is shown that the **minimal solution** of this condition reads

$$S_{nat} = \left[SU(2)_{G_L} \times SU(2)_{G_R} \times U(1)^{B-L}_{G_B}\right]_{elem} \times \left[SU(2)_{\Gamma_L} \times SU(2)_{\Gamma_R}\right]_{comp} \ . \tag{2.3}$$

The Goldstone boson matrix $\Sigma(x) \in SU(2)$ (needed to give masses to W and Z) transforms according to a different local chiral symmetry

$$\Sigma(x) \to \Gamma_L(x)\Sigma(x)\left[\Gamma_R(x)\right]^{-1} \tag{2.4}$$

than the chiral fermion doublets and the elementary gauge fields coupled to fermions

$$\psi_{L/R} \to G_{L/R} \exp\left[-i\frac{B-L}{2}\alpha\right]\psi_{L/R} \ . \tag{2.5}$$

The most general Lagrangian of dimension $d = 2$ invariant under the linear action of the symmetry S_{nat} reads

$$\begin{aligned}\mathscr{L}\left(p^2\right) \ = \ & \frac{F_W^2}{4}\left\langle D_\mu \Sigma^\dagger D^\mu \Sigma\right\rangle + i\,\overline{\psi_L}\gamma^\mu D_\mu \psi_L + i\,\overline{\psi_R}\gamma^\mu D_\mu \psi_R \\ & -\frac{1}{2}\left\langle G_{L\mu\nu}G_L^{\mu\nu} + G_{R\mu\nu}G_R^{\mu\nu}\right\rangle - \frac{1}{4}G_{B\mu\nu}G_B^{\mu\nu} \ .\end{aligned} \tag{2.6}$$

It contains several gauge fields not observed at low energies $E < \Lambda_W$, no fermion masses and a gauge boson mass term which has no obvious connection with the SM. Nevertheless, the above Lagrangian reduces to the one of the SM upon imposing S_{nat} - invariant constraints eliminating the redundant gauge fields through pairwise identification of different gauge factors up to a gauge transformation. (Notice that these constraints break the accidental L-R symmetry present in (2.6).) Example of such constraints is

$$\Gamma_{L,\mu} = \mathscr{X} g_L G_{L,\mu} \mathscr{X}^{-1} + i\mathscr{X}\partial_\mu \mathscr{X}^{-1} \tag{2.7}$$

which replaces $SU(2)_{G_L} \times SU(2)_{\Gamma_L}$ by its diagonal subgroup (identified with the SM weak isospin) and a scalar object $\mathscr{X}$ which is a (constant) multiple of a $SU(2)$ matrix, and is called "spurion". Similarly, one identifies up to a gauge $\Gamma_{R,\mu} \sim g_R G_{R,\mu} \sim g_B G_{B,\mu}\tau_3/2$. We then remain with the gauge fields of the SM, receiving standard masses and mixing through the first term in (2.6) and coupled in the standard way to fermions. In addition, we now have three $SU(2)$ valued spurions $\mathscr{X}$, $\mathscr{Y}$ and ω

$$\mathscr{X}(x) = \xi\,\Omega_L(x), \quad \Omega_L(x) \in SU(2), \quad \mathscr{Y} = \eta\,\Omega_R, \quad \Omega_R \in SU(2), \quad \omega = \zeta\,\Omega_B, \quad \Omega_B \in SU(2)\ , \tag{2.8}$$

populating the coset space $S_{nat}/S_{EW} = SU(2)^3$. To maintain invariance under S_{nat}, the spurions have to transform as

$$\mathscr{X} \to \Gamma_L \mathscr{X} G_L^{-1}, \quad \mathscr{Y} \to \Gamma_R \mathscr{Y} G_R^{-1}, \quad \omega \to \Gamma_R\,\omega\, G_B^{-1}. \tag{2.9}$$

Consequently, the constraints selecting $S_{EW} = SU(2)_W \times U(1)_Y$ of the SM as the maximal subgroup of S_{nat} that is linearly realized at low energies can be equivalently written as

$$D_\mu \mathscr{X} = 0, \quad D_\mu \mathscr{Y} = 0, \quad D_\mu \omega = 0 \tag{2.10}$$

indicating that spurions do not propagate. There exists a gauge in which the spurions reduce to three real parameters ξ, η and ζ which are exterior to the SM and whose magnitude is not fixed by the LEET. They will be considered as **small expansion parameters** describing effects beyond the SM.

The physical origin of spurions satisfying the constraints (2.10) can be understood as resulting from a particular non decoupling limit of an ordinary Higgs mechanism in which both Higgs bosons and some combinations of gauge fields become very massive. Massive gauge fields decouple, whereas heavy Higgs fields reduce to non propagating spurions, defining a non linear realization of the symmetry S_{nat}/S_{EW}.

Spurions are **needed** to write down S_{nat} invariant fermion masses. Consequently, the latter will be suppressed with respect to the scale Λ_W by powers of spurion parameters ξ and η. The least suppressed mass - the top mass - will be proportional to the product

$$\xi\eta \sim m_{top}/\Lambda_W = \mathscr{O}(p), \quad d^* = d + \frac{1}{2}(n_\xi + n_\eta). \tag{2.11}$$

This suggests to extend the low-energy power counting to spurions introducing the chiral dimension $d^\star$ defined above. This guarantees that both the fermion mass term and the lagrangian (2.6) have $d^\star = 2$ characteristic of the leading order of the LEET. Notice that the power counting formula also holds replacing in (2.2) d by $d^\star$.

The third spurion ω breaks B-L, which is thus predicted to be a part of the LEET. Consequently, the parameter $\zeta \ll \xi \sim \eta$ naturally accommodates the smallness of Lepton number violation and of the Majorana masses.

3. Next to Leading Order (NLO)

The NLO consists of all S_{nat} invariant operators of the chiral dimension $d^\star = 3$. There are two and only two such operators: they describe non standard couplings of fermions to W and Z and they are suppressed by two powers of spurions $\mathscr{X}$ or $\mathscr{Y}$:

$$\mathscr{O}_L = \bar{\psi}_L \mathscr{X}^\dagger \gamma^\mu \Sigma D_\mu \Sigma^\dagger \mathscr{X} \psi_L, \tag{3.1}$$

for left handed fermions, whereas for right handed fermions one has

$$\mathscr{O}_R^{a,b} = \bar{\psi}_R \mathscr{Y}_a^\dagger \gamma^\mu \Sigma^\dagger D_\mu \Sigma \mathscr{Y}_b \psi_R. \tag{3.2}$$

where $a,b \in [U,D]$, label covariant projections on Up and Down components of right handed doublets. These operators already carry their respective suppression factors, they are $\mathscr{O}(p^2\xi^2)$ and $\mathscr{O}(p^2\eta^2)$ respectively. The full $d^\star = 3$ part of the effective Lagrangian can be written as

$$\mathscr{L}_{\mathrm{NLO}} = \rho_L \mathscr{O}_L(l) + \lambda_L \mathscr{O}_L(q) + \sum_{a,b} \rho_R^{a,b} \mathscr{O}_R^{a,b}(l) + \sum_{a,b} \lambda_R^{a,b} \mathscr{O}_R^{a,b}(q) \tag{3.3}$$

where ρ and λ are dimensionless low-energy constants which should be of order one (unless suppressed by an additional symmetry). The NLO couplings still respect the family symmetry . On the

other hand, at this subleading order, the lepton - quark universality could be broken, i.e. $\rho \neq \lambda$ by the existence of additional reflection symmetry $\nu_R \to -\nu_R$ which does not exist for quarks. Such a symmetry is not obstructed by the LO couplings to gauge fields (at LO, ν_R decouples). It allows the right handed neutrino to get a **small Majorana mass** of the order $\mathscr{O}(\zeta^2\eta^2)$, i.e. of a comparable size to left handed Majorana mass $\mathscr{O}(\zeta^2\xi^2)$ and to the strength of LNV. On the other hand, the reflection symmetry $\nu_R \to -\nu_R$ forbids the Dirac neutrino-mass and could provide a natural explanation of the observed smallness of neutrino masses. A corollary of this "anti see-saw" mechanism [3] of suppression of neutrino masses is the **suppression of charged leptonic right-handed currents** i.e. $\rho_R^{UD} = 0$ in Eq (3.3).

4. Couplings to W

Let us concentrate on couplings of fermions to W. Using the matrix notation in the family space $\mathrm{U} = (u,c,t)^T$, $\mathrm{D} = (d,s,b)^T$, $\mathrm{N} = (\nu_e,\nu_\mu,\nu_\tau)^T$, $\mathrm{L} = (e,\mu,\tau)^T$ and using the mass - diagonal basis, the couplings to W up to and including NLO become

$$\mathscr{L}_{\mathrm{W}} = \frac{e(1-\xi^2\rho_L)}{\sqrt{2}s}\left\{\bar{\mathrm{N}}_L V_{\mathrm{MNS}}\gamma^\mu L_L + (1+\delta)\bar{\mathrm{U}}_L V_L \gamma^\mu D_L + \varepsilon \bar{\mathrm{U}}_R V_R \gamma^\mu D_R\right\} W_\mu^+ + \mathrm{h.c}\,. \tag{4.1}$$

V_L and V_R are two independent **unitary** mixing matrices resulting (as in SM) from the diagonalization of quark masses. The (small) spurionic parameters $\delta = (\rho_L - \lambda_L)\xi^2$ and $\varepsilon = \lambda_R^{UD}\eta^2$ describe the chiral generalization of the CKM mixing induced by RHCs. Notice, in particular, that effective EW couplings in the vector and axial channels (more directly accessible than V_L and V_R)

$$\begin{aligned} \mathscr{V}_{\mathit{eff}}^{ij} &= (1+\delta)V_L^{ij} + \varepsilon V_R^{ij} + \mathrm{NNLO}\,, \\ \mathscr{A}_{\mathit{eff}}^{ij} &= -(1+\delta)V_L^{ij} + \varepsilon V_R^{ij} + \mathrm{NNLO}\,, \end{aligned} \tag{4.2}$$

need not to be unitary. The signal of RHCs can be detected as $\mathscr{V}_{\mathit{eff}}^{ij} \neq -\mathscr{A}_{\mathit{eff}}^{ij}$, i.e. comparing vector and axial vector transitions.

A particular attention should be payed to **light quarks** u,d,s for which the chirality breaking effects are tiny. In this sector all EW effective couplings can be expressed in terms of δ and three parameters

$$\varepsilon_{ns} = \varepsilon\ \mathrm{Re}\Big(\frac{V_R^{ud}}{V_L^{ud}}\Big), \quad \varepsilon_s = \varepsilon\ \mathrm{Re}\Big(\frac{V_R^{us}}{V_L^{us}}\Big)\,, \quad \mathscr{V}_{\mathit{eff}}^{ud} = 0.97377(26) \equiv cos\hat{\theta} \tag{4.3}$$

where $\mathscr{V}_{\mathit{eff}}^{ud}$ is determined from $0^+ \to 0^+$ nuclear transitions [5]. Using further the unitarity of V_L and neglecting $|V_L^{ub}|^2$, all light quark effective couplings can be expressed as

$$\begin{aligned} |\mathscr{V}_{\mathit{eff}}^{ud}|^2 &= \cos^2\hat{\theta} \\ |\mathscr{A}_{\mathit{eff}}^{ud}|^2 &= \cos^2\hat{\theta}\,(1-4\,\varepsilon_{ns}) \\ |\mathscr{V}_{\mathit{eff}}^{us}|^2 &= \sin^2\hat{\theta}\,(1+2\frac{\delta+\varepsilon_{ns}}{\sin^2\hat{\theta}})(1+2\,\varepsilon_s - 2\,\varepsilon_{ns}) \\ |\mathscr{A}_{\mathit{eff}}^{us}|^2 &= \sin^2\hat{\theta}\,(1+2\frac{\delta+\varepsilon_{ns}}{\sin^2\hat{\theta}})(1-2\,\varepsilon_s - 2\,\varepsilon_{ns})\,. \end{aligned} \tag{4.4}$$

The genuine spurion parameters δ and ε are expected to be at most of order few percent. Since $|V_L^{us}| \ll |V_L^{ud}| \sim 1$ and the matrix V_R is unitary, one should have $|\varepsilon_{NS}| < \varepsilon$. On the other hand, the parameter ε_S measuring RHCs strangeness changing transitions **can be enhanced** if the mixing hierarchy for right handed light quarks is inverted, $V_R^{ud} < V_R^{us}$. In this case $|\varepsilon_S|$ could be as large as 4.5ε. Clearly, this question should be decided experimentally.

5. The stringent test of RHCs: Scalar $K_{\mu 3}$ form factor shape

Model independent bounds on $V+A$ couplings of light quarks to W are extremely difficult to find, since they require an accurate control of QCD chiral symmetry breaking contributions when comparing hadronic matrix elements of vector and axial vector currents. One such test (never considered before) has been identified in Ref [7]. It is based on the Callan Treiman low - energy Theorem already discussed in the talk by E. Passemar [6]. The (normalized) scalar $K_{\mu 3}^L$ form factor $f(t)$

$$f(t) = \frac{f_S^{K^0\pi^-}(t)}{f_+^{K^0\pi^-}(0)} = \frac{1}{f_+^{K^0\pi^-}(0)}\left(f_+^{K^0\pi^-}(t) + \frac{t}{\Delta_{K\pi}} f_-^{K^0\pi^-}(t)\right) \;,\quad f(0)=1. \tag{5.1}$$

where $\Delta_{K\pi} = m_K^2 - m_\pi^2$, satisfies

$$C \equiv f(\Delta_{K\pi}) = \frac{F_{K^+}}{F_{\pi^+}} \frac{1}{f_+^{K^0\pi^-}(0)} + \Delta_{CT}, \quad C = B_{exp}\, r + \Delta_{CT}. \tag{5.2}$$

Here, $\Delta_{CT} = -3.5\times 10^{-3}$ is a tiny correction which has been estimated in one loop ChPT. In the absence of RHCs, the value C of the scalar form factor at the Callan Treiman point can be directly expressed in terms of measured branching fractions (K_{l2}/π_{l2}, K_{l3}) and V^{ud} giving [5] $B_{exp} = 1.2438 \pm 0.0040$ in the second Eq. (5.2). RHCs make appear additional correction factor r

$$r = \left| \frac{\mathscr{A}_{eff}^{ud}\mathscr{V}_{eff}^{us}}{\mathscr{V}_{eff}^{ud}\mathscr{A}_{eff}^{us}} \right| = 1 + 2(\varepsilon_S - \varepsilon_{NS}). \tag{5.3}$$

Hence, in the presence of RHCs the Callan Treiman theorem yields

$$\ln C = 0.2182 \pm 0.0035 + \tilde{\Delta}_{CT} + 2(\varepsilon_s - \varepsilon_{ns}) = 0.2182 \pm 0.0035 + \Delta\varepsilon \;, \tag{5.4}$$

with $\tilde{\Delta}_{CT} = \Delta_{CT}/B_{exp}$.

An accurate physically motivated parametrization of the scalar form-factor $f(t)$ has been proposed [7] which allows to determine the parameter $\ln C$ from the measured $K_{\mu 3}^L$ decay distributions. The corresponding measurement is particularly delicate, since the experimental t - distribution is not easy to reconstruct from the data. Furthermore, different experiments have access to different decay distributions which do not have the same sensitivity to $\ln C$ and to the shape of the vector form factor. There exists a relation between $\ln C$ and the slope parameter λ_0 [6] but it is not enough precise to reduce the determination of $\ln C$ to existing (controversial) determinations of the slope λ_0 assuming the linear t-dependence of the scalar form factor [8, 9, 10] or at most injecting information about its curvature [11]. Recently, NA48 collaboration has published the result of a direct determination of $\ln C$ based on the dispersive representation of $f(t)$ [10]

$$\ln C_{exp} = 0.1438 \pm 0.0138, \quad \Delta\varepsilon = -0.074 \pm 0.014 \tag{5.5}$$

Other analysis of $K_{\mu 3}$ decay distributions from KLOE [12] and KTeV [13] based on the dispersive representation of the two form factors are underways. They should clarify the experimental situation and provide an independent cross check of the NA48 result [10]. Awaiting an independent dispersive analysis of existing data samples, one should stress that the result (5.5) indicates a 5σ deviation from the SM prediction. In particular, if the discrepancy would have to be explained within QCD, the ChPT estimate of Δ_{CT} would have to be underestimated by a factor 20. On the other hand, within the class of LEET defined above the interpretation of the result (5.5) as a manifestation of couplings of right handed quarks to W is unambiguous. It amounts to a determination of the spurion parameter $2(\varepsilon_S - \varepsilon_{NS})$. Its size can be understood as a result of enhancement of V_R^{us} relative to the suppressed V_L^{us}. Beyond our LEET framework, other interpretations might be conceivable. For example a subTeV charged scalar coupled to scalar densities $\bar{u}s$ and $\bar{\mu}\nu$ could interfere with our analysis. We prefer to stay within the class of minimal LEET defined above and ask how does the same non standard operator (3.2) affect the couplings of right handed quarks to Z.

6. Couplings to Z

Non standard couplings to Z contained in the NLO Lagrangian (3.3) are suppressed by the same two spurion parameters ξ^2 (LHCs) and η^2 (RHCs) as in the case of couplings to W discussed in Section 4. Hence, despite the apriori unknown "order one" prefactors ρ and λ, it is possible to relate orders of magnitude of non standard CC and NC couplings. In the left - handed sector we have altogether two NLO parameters: $\delta = \xi^2(\rho_L - \lambda_L)$ and $\xi^2 \rho_L$, whereas in the right - handed sector there are three new parameters denoted ε^e, ε^U, ε^D and proportional to the spurion η^2.

	Measurement	Fit
Γ_Z [GeV]	2.4952(23)	2.4943
σ_{had} [nb]	41.540(37)	41.569
R_e	20.767(25)	20.785
A^l_{FB}	0.0171(10)	0.0165
$\mathcal{A}_l(P_\tau)$	0.1465(32)	0.1485
R_b	0.21629(66)	0.21685
R_c	0.1721(30)	0.1725
A^b_{FB}	0.0992(16)	0.1012
A^c_{FB}	0.0707(35)	0.0707
$\mathcal{A}_b$	0.923(20)	0.910
$\mathcal{A}_c$	0.670(27)	0.636
$\mathcal{A}_l$(SLD)	0.1513(21)	0.1485
$Br(W \to l\nu)$	0.1084(9)	0.1089

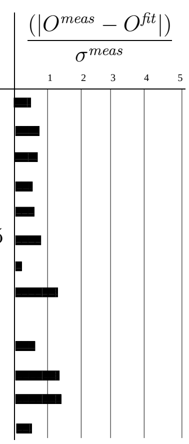

Figure 1: *Pull for the Z pole observables in the full fit*

We have performed the NLO fit to the usual set of Z - pole pseudo-observables displayed in Fig. 1 including the lepton branching fraction of W (particularly sensitive to the parameter δ) as well as spin asymmetries measured at SLD. The fit is described in details in [4]. It has $\chi^2/dof = 8.5/8$ and it gives $\delta \equiv \xi^2(\rho_L - \lambda_L) = -0.004(2)$, $\xi^2 \rho_L = 0.001(12)$ and $\varepsilon^e \equiv \eta^2 \rho_R^{DD} = -0.0024(5)$.

The most important NLO modification of couplings to Z turns out to occur for right handed quarks: $\varepsilon^U \equiv \eta^2\lambda_R^{UU} = -0.02(1) \quad \varepsilon^D \equiv \eta^2\lambda_R^{DD} = -0.03(1)$. The correlations can be found in [4].

Two comments are in order. First, the most important NLO non standard couplings to Z seem to occur for **right handed quarks**. Their size compares well with the couplings of right handed quarks to W as suggested by the $K^L_{\mu 3}$ dispersive Dalitz plot analysis [10]. Next, the fit is of a very good quality as illustrated in Fig. 1 in terms of "pulls". In particular, the b-quark forward backward asymmetry A^b_{FB} **and** R^b are both well reproduced without modifying the flavour universality of NC EW couplings. The long standing "puzzle of b-asymmetries" has apparently gone thanks to the modified right-handed couplings of D type quarks to Z.

7. F_K/F_π and $f_+(0)$

The low-energy QCD quantities $F_K, F_\pi, f_+(0)\ldots$ are defined independently of EW interactions in terms of QCD correlation functions and they are accessible to ChPT and lattice studies. On the other hand their experimental values extracted from semileptonic branching fractions depend on the presumed EW vertices via the effective EW couplings (4.4). Fixing experimental values of $\mathscr{V}^{ud}_{\mathit{eff}}$ (4.3) and of the semi leptonic branching ratios, $F_K, F_\pi, f_+(0)\ldots$ become unique functions of spurion parameters ε_{NS}, ε_S and δ. One has

$$\left(\frac{F_{K^+}}{F_{\pi^+}}\right)^2 = \left(\frac{\hat{F}_{K^+}}{\hat{F}_{\pi^+}}\right)^2 \frac{1+2\,(\varepsilon_s-\varepsilon_{ns})}{1+\frac{2}{\sin^2\hat{\theta}}(\delta+\varepsilon_{ns})}, \quad |f_+^{K^0\pi^-}(0)|^2 = \left[\hat{f}_+^{K^0\pi^-}(0)\right]^2 \frac{1-2(\varepsilon_s-\varepsilon_{ns})}{1+\frac{2}{\sin^2\hat{\theta}}(\delta+\varepsilon_{ns})}, \tag{7.1}$$

where the hat indicates the corresponding values extracted from semi leptonic branching fractions (**assuming SM couplings** $\varepsilon_{NS} = \varepsilon_S = \delta = 0$). The latter are known very precisely:

$$\hat{F}_{K^+}/\hat{F}_{\pi^+} = 1.182(7), \quad \hat{f}_+^{K^0\pi^-}(0) = 0.951(5)\ . \tag{7.2}$$

In fig. 2 are displayed lines of constant values of F_K/F_π and $f_+(0)$ as a function of spurion parameters. One notes that F_K/F_π significantly decreases compared with the value 1.22 often used as input in ChPT. On the other hand, $f_+(0)$ is not very constrained despite the Callan Treiman relation. Finally, nothing prevents the effective vector mixing matrix $\mathscr{V}_{eff}$ to be nearly unitary without any fine tuning. One has

$$|\mathscr{V}^{ud}_{\mathit{eff}}|^2 + |\mathscr{V}^{us}_{\mathit{eff}}|^2 = 1 + 2\,(\delta+\varepsilon_{ns}) + 2\,(\varepsilon_s-\varepsilon_{ns})\,\sin^2\hat{\theta}\ . \tag{7.3}$$

The contribution of ε_S, the only parameter which might be enhanced above 0.01 is suppressed by $sin^2\hat{\theta}$.

In conclusion, we have presented and motivated a new low-energy test of non standard EW couplings of right handed quarks not considered before. Did one observe couplings of right handed quarks to W in $K^L_{\mu 3}$ decay ? The final answer requires a more complete and dedicated experimental analysis. It also deserves a particular effort despite its difficulty.

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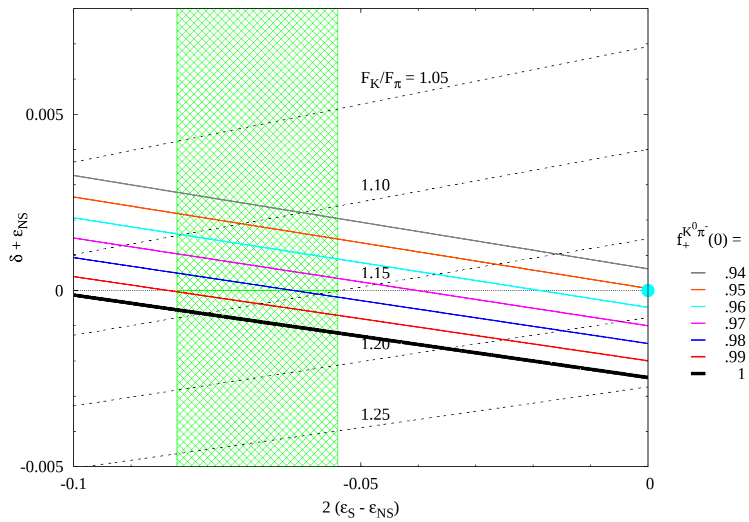

Figure 2: *Lines of constant values for $F_{K+}/F_{\pi+}$ and $f_+^{K^0\pi}(0)$ in the plane $\delta+\varepsilon_{ns}$ and $2(\varepsilon_s-\varepsilon_{ns})$ as resulting from Eqs (7.1) and (7.2). The vertical band indicates the range suggested by the NA48 measurement [10]. The SM point $\varepsilon=\delta=0$ is also shown.*

References

[1] R.N. Mohapatra and J.C. Pati, *Phys. Rev. D* **11** (1975) 2558; G. Senjanovic and R. N. Mohapatra, *Phys. Rev. D* **12** (1975) 1502

[2] J. Hirn and J. Stern, *Eur. Phys. J. C* **34** (2004) 447 [`hep-ph/0401032`]; J. Hirn and J. Stern, *JHEP* **0409** (2004) 058 [`hep-ph/0403017`]

[3] J. Hirn and J. Stern, *Phys. Rev. D* **73** (2006) 056001 [`hep-ph/0504277`]

[4] V. Bernard, M. Oertel, E. Passemar and J. Stern, [`hep-ph/0707.4194`]

[5] E. Blucher and W. Marciano in W. M. Yao *et al.* [Particle Data Group], *J. Phys. G* **33** (2006) 1

[6] E. Passemar, these Proceedings

[7] V. Bernard, M. Oertel, E. Passemar and J. Stern, *Phys. Lett. B* **638** (2006) 480 [`hep-ph/0603202`]

[8] M. Palutan, these Proceedings

[9] T. Alexopoulos *et al.* [KTeV Collaboration], *Phys. Rev. D* **70** (2004) 092007 [`arXiv:hep-ex/0406003`]

[10] A. Lai *et al.* [NA48 Collaboration], *Phys. Lett. B* **647** (2007) 341 [`hep-ex/0703.002`]; R. Wanke, these Proceedings

[11] F. Ambrosino *et al.* [KLOE Collaboration], [`hep-ex/0707.4631`]

[12] M. Antonelli, Private Communication

[13] A. Glazov, Private Communication